

Light inflaton after LHC8 and WMAP9 results

F. Bezrukov

*Physics Department, University of Connecticut, Storrs, CT 06269-3046, USA and
RIKEN-BNL Research Center, Brookhaven National Laboratory, Upton, NY
11973-5000, USA
E-mail: Fedor.Bezrukov@uconn.edu*

D. Gorbunov

*Institute for Nuclear Research of the Russian Academy of Sciences, Moscow 117312,
Russia and
Moscow Institute of Physics and Technology, Dolgoprudny 141700, Russia
E-mail: gorby@ms2.inr.ac.ru*

ABSTRACT: We update the allowed parameter space of the simple chaotic inflationary model with quartic potential and light inflaton [1] taking into account recent results from cosmology (CMB observations from SPT, ACT and WMAP) and from particle physics (LHC hints of the SM Higgs boson). The non-minimal (yet small) coupling to gravity of the inflaton becomes essential to fit the observational data. The inflaton has mass above 300 MeV and can be searched for at B -factories in B -meson two-body decays to kaon and inflaton. The inflaton lifetime depends on the model parameters, resulting in various inflaton signatures: either a missing energy, or a displaced vertex from the B -meson decay position, or a resonance in the Dalitz plot of a three particle decay. We also discuss the implementation of the inflaton model to the ν MSM, where the inflaton can be responsible for production of the dark matter sterile neutrino in the early Universe.

KEYWORDS: [inflation](#), [particle physics](#).

Contents

1. Introduction	1
2. Model: inflaton with $\beta X^4/4$ potential, non-minimal coupling to gravity and mixing with the SM Higgs boson	3
3. Inflation and reheating: prediction of cosmological parameters	4
3.1 Slow roll parameters and spectral indices	4
3.2 Horizon crossing and reheating	7
4. Phenomenology: searches in particle physics experiments	10
5. Discussion	13

1. Introduction

Elegant solutions to the major problems of the Hot Big Bang theory are provided by a preliminary inflationary stage of the Universe expansion [2, 3, 4, 5, 6]. The simplest realization is made using a scalar degree of freedom, inflaton, which evolves very slowly at the inflationary stage providing nearly constant energy density and hence the accelerated expansion. At later times the inflaton starts to oscillate and the Universe exits inflation and enters a stage when rapidly evolving homogeneous inflaton dominates the total energy density. The energy has to be transferred to the Standard Model (SM) particles thus reheating the Universe. To make the transfer an inflaton-to-SM coupling is needed. This coupling takes care of the SM particle production and the onset of Hot stage of the Universe expansion.

Let us analyse the consequences of adding explicitly just one scalar inflaton field to the SM. In a renormalizable model of the inflaton sector the coupling to the SM should be renormalizable as well. Then for the simplest inflationary models with only one scalar field, inflaton, involved in the inflationary dynamics, the only suitable coupling to the SM is its coupling to the SM Higgs field via the scalar potential. Then reheating happens via the Higgs boson production by oscillating inflaton field. A variant of these models has been considered in Refs. [7, 1]. It is distinguished by the absence of dimensionful parameters in the sector of SM fields. The vacuum expectation value of the SM Higgs field v appears as a nonzero vacuum expectation value of the inflaton field: the quartic coupling between the two field is responsible for both reheating in the post-inflationary Universe and electroweak phase transition later at the Hot stage. The only dimensionful parameter is in the inflaton sector, and it is of the order of electroweak scale. All model parameters are therefore stable with respect to the finite quantum corrections (or logarithmic renormalization running).

The appealing phenomenological feature of the inflationary model described above is the new light particle, inflaton, with the light mass in GeV region. Cosmology limits the mass range of the light inflaton from above and from below: lighter inflaton causes very late reheating, while quantum corrections from the inflaton-to-Higgs quartic coupling spoil flatness of the inflaton potential for heavy inflaton masses. The latter bound on the inflaton-to-SM coupling constrains the inflaton mass to be below 2 GeV [1].

In the model the quartic coupling induces the Higgs-inflaton mixing, so the inflaton mass is proportional to the mass of the Higgs boson and square root of the quartic coupling. The inflaton mass range presented above refers to the Higgs mass of 100–200 GeV. Recent results from LHC [8, 9] can be interpreted as evidence for the SM 125 GeV Higgs boson, which allows to resolve the ambiguity due to the Higgs mass value: now the only free parameter remaining in the model is the Higgs-inflaton quartic coupling, or equivalently the inflaton mass. Thus the estimates given in Ref. [1] can be appropriately refined. It is important for the phenomenology of the light inflaton, studied there. Indeed, provided mixing with the SM Higgs boson, inflaton may be produced in scatterings and decays of the SM particles and subsequently decays into the SM particles. The latter branching ratios are the same as those for a (hypothetical) SM Higgs boson of mass equal to those of the inflaton. Thus, *the inflaton mass alone determines the inflaton production rate in the laboratory experiment and its lifetime*, which was estimated as 10^{-8} – 10^{-10} s. The most promising place to search for the inflaton was suggested to be in B -meson decays [1].

Another issue has risen due to the recent analyses of cosmological data, mainly observations of the Cosmic Microwave Background (CMB) anisotropy by SPT [10], ACT [11], and the common global fit together with the analysis of CMB data collected for 9 years by WMAP [12]. The most important, relevant for our model result is lowering an upper limit on the tensor-to-scalar perturbations ratio, which excludes the simple chaotic inflation with quartic potential discussed above at more than 95% confidence level. It was already pointed in Ref. [1] that with a small non-minimal coupling to gravity the situation may be cured. It was found [1] that, at rather small value of non-minimal coupling constant $\lesssim 10^{-3}$, this also leads to the increase of the inflaton self-coupling constant. The latter is determined from the amplitude of the scalar perturbation power spectrum obtained mostly from analyses of CMB anisotropy. This change was accounted for in the estimates presented in Ref. [1]. However, with new results from CMB anisotropy [10, 12, 11] a somewhat larger non-minimal coupling is needed to make the inflaton model viable, compared to the value adopted in Ref. [1]. Hence that analysis has to be modified appropriately.

In the present paper we update the analysis of Ref. [1] to account for both recent measurement of the Higgs boson mass of about 126 GeV [8, 9] and recent new upper limit on the tensor-to-scalar ratio $r < 0.13$ (at 95% C.L. [12]).

The paper is organized as follows. In Sec. 2 we present the model Lagrangian and discuss the particle spectrum in the scalar sector. In Section 3, which is devoted to inflationary dynamics of the model, we give predictions for cosmological parameters to be tested in future cosmological experiments. Particle physics phenomenology is discussed in Sec. 4, where we refine the estimates of the mass range of the light inflaton and discuss the strategy of searches for light inflaton in particle physics experiments, such as LHCb and

Belle-II. Section 5 contains conclusions and further discussion on the model extension to solve the major phenomenological problems of the SM: neutrino oscillations, dark matter phenomenon and baryon asymmetry of the Universe; here an example is ν MSM [13, 14].

2. Model: inflaton with $\beta X^4/4$ potential, non-minimal coupling to gravity and mixing with the SM Higgs boson

The action of the light inflaton model is [1]

$$S_{XSM} = \int \sqrt{-g} d^4x (\mathcal{L}_{SM} + \mathcal{L}_{XN} + \mathcal{L}_{\text{ext}} + \mathcal{L}_{\text{grav}}),$$

$$\mathcal{L}_{XN} = \frac{1}{2} \partial_\mu X \partial^\mu X + \frac{1}{2} m_X^2 X^2 - \frac{\beta}{4} X^4 - \lambda \left(H^\dagger H - \frac{\alpha}{\lambda} X^2 \right)^2, \quad (2.1)$$

$$\mathcal{L}_{\text{grav}} = - \frac{M_P^2 + \xi X^2}{2} R, \quad (2.2)$$

where \mathcal{L}_{SM} is the SM Lagrangian without the Higgs field potential, and \mathcal{L}_{ext} stands for an extension of the SM capable of explaining all the major phenomenological puzzles: neutrino oscillations, dark matter phenomena, baryon asymmetry of the Universe (see discussion in Sec. 5). As far as we assume $\alpha, \beta \ll \lambda$, inflation proceeds along the flat direction of the scalar potential, where the Higgs and inflaton contributions in the last term of (2.1) cancel. In the inflaton sector of the scalar potential the sign of the quadratic term is chosen to lead to the nonzero vacuum expectation value for the inflaton field at late stages of the Universe evolution (after reheating). Then the last term in (2.1) gives rise to spontaneous breaking of electroweak (EW) symmetry, and the SM Higgs field gains non-zero vacuum expectation value as well. The two scalar excitations above the vacuum, h and χ , correspond to the SM Higgs boson and inflaton particles, but the mass basis is slightly rotated as compared to H and X , with the small mixing angle θ . Four parameters of the model, m_X , β , λ , and α , determine the Higgs field vacuum expectation value $v \approx 246$ GeV, the Higgs boson mass $m_h \approx 126$ GeV, and the inflaton mass

$$m_\chi = m_h \sqrt{\frac{\beta}{2\alpha}} = \sqrt{\frac{\beta}{\lambda \theta^2}}. \quad (2.3)$$

Thus, at a given value of β , the only free parameter in the scalar sector is the mixing coupling α or the inflaton mass m_χ , which governs the inflaton effective coupling to all the SM fields via mixing with the Higgs boson characterized by the squared mixing angle

$$\theta^2 = \frac{2\beta v^2}{m_\chi^2} = \frac{2\alpha}{\lambda}. \quad (2.4)$$

For zero or very small non-minimal coupling constant ($\xi < 10^{-4}$) the inflationary dynamics is fully determined by the parameter β , which is then fixed from the amplitude of primordial density perturbations. This amplitude is measured from observations of CMB anisotropy, and the resulting value, used in Ref. [1], is $\beta = (1 - 2) \times 10^{-13}$. In this case the only free parameter in the model is the inflaton mass m_χ (uniquely related to α for given Higgs

mass and β , see (2.4)). Further, stability of the inflaton potential with respect to quantum corrections from the SM sector due to inflaton mixing to the Higgs field places an upper limit on α , while lower limit comes from the requirement of efficient reheating to the temperature exceeding 100 GeV (to allow for some baryogenesis mechanism). Consequently, using the bounds for α , the inflaton mass was confined in [1] between 30 MeV and 2 GeV. In Sec. 3 we extend those estimates for larger ($\xi > 10^{-4}$) values of non-minimal coupling constant ξ favored from the combined analysis of present cosmological data [12].

Let us note that generally one more mass parameter in the SM Higgs sector may be expected (corresponding to the quadratic term for the Higgs field) and one more mass parameter in the inflaton sector (the one in front of the cubic term). The latter is even welcome, as it helps to get rid of domain walls emerging in the Universe at the moment when inflaton gains non-zero vacuum expectation value and Z_2 symmetry of (2.1) becomes spontaneously broken. In case the new dimensionful parameters are of the order of EW scale or below, they do not change the scale of inflaton mass, but relax the relations between inflation-related and low-energy physics parameters: the value of inflaton mass would not uniquely determine the inflaton couplings to other fields in that case. In what follows we assume those parameters are much smaller than the EW scale and hence negligible for our analysis.

3. Inflation and reheating: prediction of cosmological parameters

The inflation in the model happens along the direction $H^\dagger H = \frac{\alpha}{\lambda} X^2$, so that the potential term in (2.1), which has a not quite small coupling constant λ , vanishes.¹ The mass term in (2.1) is negligible because m_X is significantly below the inflationary scale. Thus, the only relevant terms during the inflation are the quartic potential term and non-minimal coupling with gravity. This brings us to the situation thoroughly analyzed in literature (see, e.g. [15, 16]). Below in Sec. 3.1 we repeat the analysis to arrive at already known results and obtain some new analytic approximations valid in the interesting part of the parameter space.

3.1 Slow roll parameters and spectral indices

To apply the standard slow roll formalism (see e.g. [17, 18]) we perform a conformal transformation of the metric

$$g_{\mu\nu} \rightarrow \tilde{g}_{\mu\nu} = \Omega^2 g_{\mu\nu}, \quad \Omega^2 = 1 + \xi X^2/M_P^2, \quad (3.1)$$

In the new variables $\tilde{g}_{\mu\nu}$ (Einstein frame, as opposed to the original Jordan frame) gravity couples minimally to all matter fields including inflaton. Then the inflaton potential gets modified and reads

$$U(X) = \frac{\beta X^4}{4\Omega^4}. \quad (3.2)$$

¹This is true for $\beta \ll 1$, which is the case we are interested in this article.

After the conformal transformation the inflaton kinetic term gets modified too. In the Einstein frame the canonically normalized “inflaton field” \mathcal{X} is obtained from the solution of

$$\frac{d\mathcal{X}}{dX} = \sqrt{\frac{\Omega^2 + 6\xi^2 X^2/M_P^2}{\Omega^4}}. \quad (3.3)$$

Then the potential term for the field \mathcal{X} is (3.2) where the solution $X = X(\mathcal{X})$ of eq. (3.3) is substituted.

For the canonically normalized field \mathcal{X} we can use the standard formulas for the slow roll parameters during inflation, number of e-folding and spectral indices. It has been proven for our model that prediction for cosmological parameters does not depend on the choice of frame [16]. In the Einstein frame we have a usual large field chaotic inflation with the potential (3.2), (3.3). In this case the amplitudes and spectral indices of scalar and tensor perturbations, prediction of inflation, are generally determined by the value of the scalar potential and values of the slow roll parameters. The slow roll parameters are [15]

$$\epsilon = \frac{M_P^2}{2} \left(\frac{dU/d\mathcal{X}}{U} \right)^2 = \frac{8}{(1 + \xi(1 + 6\xi)X^2/M_P^2)X^2/M_P^2}, \quad (3.4)$$

$$\eta = M_P^2 \frac{d^2 U/d\mathcal{X}^2}{U} = \frac{4(3 + \xi(1 + 12\xi)X^2/M_P^2 - 2\xi^2(1 + 6\xi)X^4/M_P^4)}{(1 + \xi(1 + 6\xi)X^2/M_P^2)^2 X^2/M_P^2}. \quad (3.5)$$

At the end of inflation $\epsilon = 1$ and then from (3.4) the field value X_e is given by [15]

$$\frac{X_e^2}{M_P^2} = \frac{\sqrt{192\xi^2 + 32\xi + 1} - 1}{2\xi(1 + 6\xi)}. \quad (3.6)$$

During inflation the inflaton quantum fluctuations become a source of matter perturbations. We are interested in the moment during inflation when perturbations of WMAP pivot scale exit the horizon due to stretching in almost exponentially expanding Universe. To estimate the inflaton field value X_N at the moment, corresponding to N e-foldings before the end of inflation, we should invert the relation [15]

$$\begin{aligned} N &= \int_{X_e}^{X_N} \frac{1}{M_P^2} \frac{U}{dU/d\mathcal{X}} \left(\frac{d\mathcal{X}}{dX} \right)^2 dX = \int_{X_e}^{X_N} \frac{X((6\xi^2 + \xi)X^2/M_P^2 + 1)}{4(1 + \xi X^2/M_P^2)M_P^2} dX \\ &= \frac{3}{4} \left[\left(\xi + \frac{1}{6} \right) (X_N^2/M_P^2 - X_e^2/M_P^2) - \ln \frac{1 + \xi X_N^2/M_P^2}{1 + \xi X_e^2/M_P^2} \right]. \end{aligned} \quad (3.7)$$

Solving this transcendental equation for X_N^2 analytically is impossible. To obtain accurate predictions we solve it numerically and present the results in Fig. 1; below we describe also the useful analytical approximations.

Now we have all formulas to match parameters of scalar and tensor perturbation spectra probed in cosmological experiments. We start with the perturbation amplitudes. One adopts the WMAP normalization of matter perturbation power spectrum [12] to fix the ratio

$$U/\epsilon = 24\pi^2 \Delta_{\mathcal{R}}^2 M_P^4 \simeq (0.0276 \times M_P)^4, \quad (3.8)$$

that in our model (3.2), (3.4) is

$$\frac{U}{\epsilon} = \frac{\beta(\xi(6\xi + 1)X_N^2/M_P^2 + 1)X_N^6/M_P^6}{32(1 + \xi X_N^2/M_P^2)^2} \quad (3.9)$$

and allows to determine the value of β at given ξ and N as

$$\beta = 24\pi^2 \Delta_{\mathcal{R}}^2 \frac{32(1 + \xi X_N^2/M_P^2)^2}{(\xi(6\xi + 1)X_N^2/M_P^2 + 1)X_N^6/M_P^6}, \quad (3.10)$$

For numerically solved (3.7) at relevant $N = 60$ (see below Sec. 3.2), we present β as function of ξ in Fig. 1, lower right panel.

Let us proceed with tensor-to-scalar ratio r and spectral indices of scalar and tensor perturbations ($n_s - 1$) and n_T , respectively. To the leading order in slow roll parameters (3.4), (3.5) the general relations are linear (see e.g. [17, 18]):

$$r = 16\epsilon, \quad n_s - 1 = 2\eta - 6\epsilon, \quad n_T = -2\epsilon. \quad (3.11)$$

They must be evaluated at the horizon crossing, when $X = X_N$, see eq. (3.7).

Two opposite cases are easy to investigate: models with small and large values of ξ . Indeed, in the limit $\xi \rightarrow 0$ we are back to the chaotic inflation with quartic scalar potential [19]. Then from eqs. (3.6), (3.7) $X_N^2/M_P^2 \rightarrow 8(N + 1)$ and further from eqs. (3.4), (3.5):

$$\epsilon = \frac{1}{N + 1}, \quad \eta = \frac{3}{2(N + 1)}. \quad (3.12)$$

Plugging these results into eq. (3.11) one arrives at

$$r = \frac{16}{N + 1}, \quad n_s - 1 = -\frac{3}{N + 1}, \quad n_T = -\frac{2}{N + 1}. \quad (3.13)$$

Eq. (3.10) gives for the quartic self-coupling

$$\beta = \frac{3\pi^2 \Delta_{\mathcal{R}}^2}{2(N + 1)^3}, \quad (3.14)$$

which equals 1.5×10^{-13} for the relevant value $N=60$ (see Sec. 3.2). It was adopted in Ref. [1] as a reference value for all numerical estimates.

In the opposite limit $\xi \rightarrow \infty$ one obtains from eqs. (3.6), (3.7) at large N (i.e. neglecting logarithmic correction in (3.7)): $\xi X_N^2/M_P^2 \rightarrow 4N/3$ and further from eqs. (3.4), (3.5):

$$\epsilon = \frac{3}{4N^2}, \quad \eta = -\frac{1}{N}. \quad (3.15)$$

Plugging these results into eq. (3.11) one arrives at [20, 16]

$$r = \frac{12}{N^2}, \quad n_s - 1 = -\frac{2}{N}, \quad n_T = -\frac{3}{2N^2}. \quad (3.16)$$

Self-coupling depends on ξ as follows from eq. (3.10):

$$\beta = \frac{72\pi^2 \Delta_{\mathcal{R}}^2}{N^2} \xi^2, \quad (3.17)$$

that is $\beta \approx 4.5 \times 10^{-10} \xi^2$ for relevant $N = 60$. Thus even $\beta \sim 1$ is allowed at sufficiently large ξ , and hence the SM Higgs boson non-minimally coupled to gravity may play a role of inflaton [21]. Note that in the limit of large ξ formulas (3.16) coincide with those obtained within R^2 -inflation [2, 3, 22]; however, numerical predictions generically differ because of different reheating temperature resulting in different values of N in different models, see discussion in Sec. 3.2. In particular, the numerical predictions of R^2 -inflation [3, 22, 23] and inflation driven by the Higgs boson [24] differ [25, 26].

To find a reasonably good analytical approximation for the intermediate case of small but finite value of ξ we utilize the following relation

$$(1 + 6\xi) X_N^2 / M_P^2 = 8(N + 1), \quad (3.18)$$

that follows from eqs. (3.6), (3.7) when linear in ξ corrections are included. Quite remarkably, though formally obtained *at small* ξ , eq. (3.18) interpolates smoothly between “ $\xi \rightarrow 0$ ” and “ $\xi \rightarrow \infty$, large N ” regimes considered above, with only $\ln N/N \simeq 5\%$ relative error in X_N^2 in the large ξ limit. Putting (3.18) into (3.4), (3.5), and (3.11) one obtains formulas for tensor-to-scalar ratio r and spectral indices, interpolating between (3.13) and (3.16), cf. [27]:

$$r = \frac{16(1 + 6\xi)}{(N + 1)(1 + 8(N + 1)\xi)}, \quad (3.19)$$

$$n_s - 1 = -\frac{3(1 + 6\xi) + 8(N + 1)(5 + 24\xi)\xi + 128(N + 1)^2\xi^2}{(N + 1)(1 + 8(N + 1)\xi)^2}, \quad (3.20)$$

$$n_T = -\frac{2(1 + 6\xi)}{(N + 1)(1 + 8(N + 1)\xi)}. \quad (3.21)$$

Similarly one obtains from (3.18) and (3.10) the formula for inflaton quartic coupling,

$$\beta = \frac{3\pi^2 \Delta_{\mathcal{R}}^2}{2} \frac{(1 + 6\xi)(1 + 6\xi + 8(N + 1)\xi)}{(1 + 8(N + 1)\xi)(N + 1)^3}, \quad (3.22)$$

interpolating between (3.14) and (3.17).

The approximations (3.19)–(3.22) are functions of ξ and N . In our model with natural reheating through the Higgs boson production we obtain $N = 60$, see Sec. 3.2. In Fig. 1, we present $n_s - 1$, r and β as *numerical solutions*. The approximate formulas (3.20), (3.19), and (3.22), have relative errors of not more than 8%, 4%, and 8%, respectively, in the entire interval of ξ (and are much more precise at low ξ). From Fig. 1 one concludes that the cosmological parameters effectively reach asymptotes of zero and infinite ξ at $\xi \lesssim 10^{-4}$ and $\xi \gtrsim 0.1$, respectively.

Note finally, that the analysis in this section can be spoiled by the radiative corrections from the interaction with the Higgs boson, which gives the contribution of the order α^2 to the inflaton quartic coupling. We will somewhat arbitrarily use the bound of $\alpha^2 < 0.1\beta$ for the radiative corrections to the inflaton.

3.2 Horizon crossing and reheating

Reheating in the model happens via the SM Higgs boson production through inflaton-to-Higgs coupling in (2.1). Usually the inflationary predictions are plotted depending on the

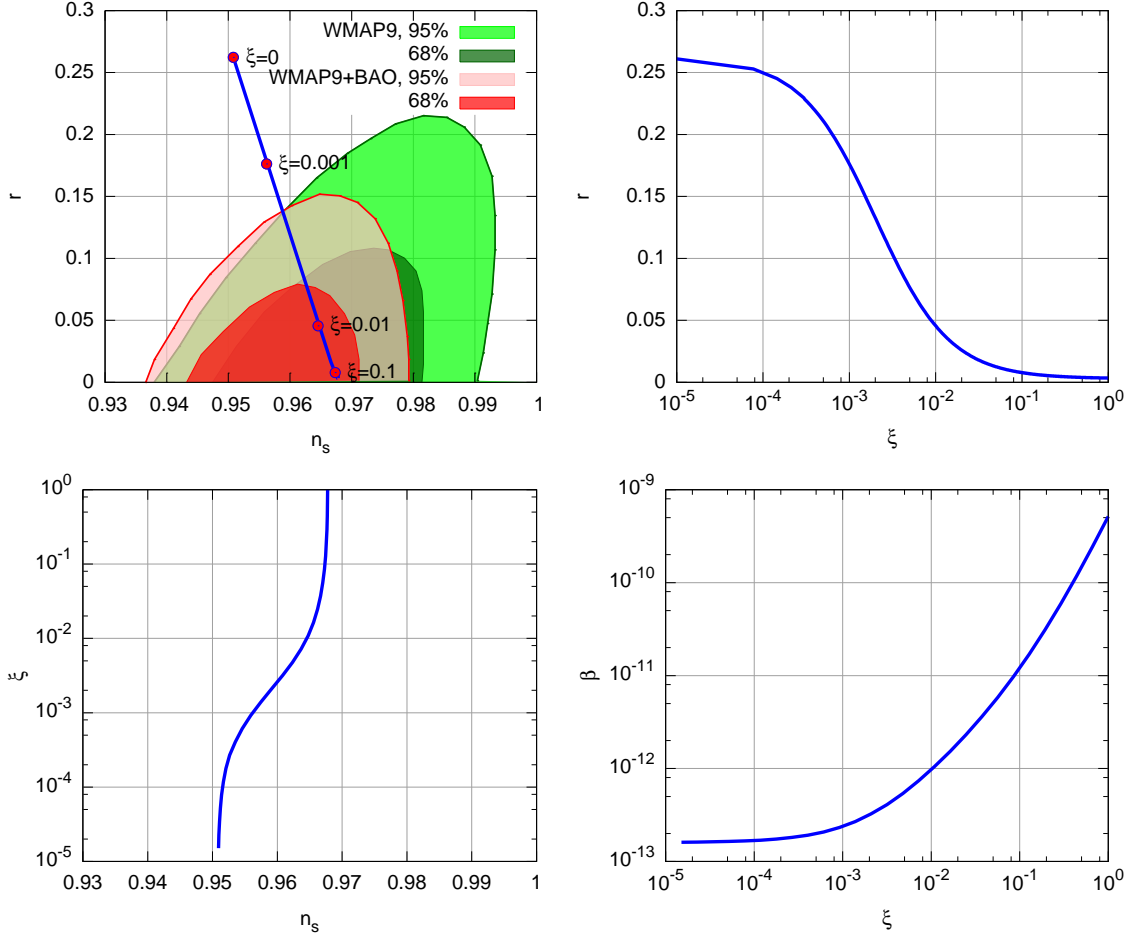


Figure 1: The cosmological parameters for $0 \leq \xi < 1$; number of e-foldings is set to $N = 60$ (see discussion in Sec. 3.2). It is clearly seen that already at $\xi = 0.1$ the model reaches the asymptote of “large” non-minimal coupling. The shaded regions on the (r, n_s) plot are outlined by 1 σ - and 2 σ -bounds from WMAP9 data (green) and WMAP9+BAO data (red) [12]. The interval $\xi < 2 \times 10^{-3}$ is disfavored from cosmology.

number of e-foldings N that happened after inflaton perturbations of a specific observable scale exit the horizon. This number is commonly estimated as $N = 50\text{--}60$. Let us stress, that in the model of this article the dynamics of the Universe expansion is known for the whole time after inflation, including reheating. Specifically, if the non-minimal coupling ξ is not too large, post-inflationary dynamics is dominated by the quartic scalar potential. This means that the Universe expands as at the radiation dominated stage. The actual reheating moment, though happening quite late [27], leads only to redistribution of energy between different relativistic degrees of freedom, and does not influence the rate of the Universe expansion.

Let us derive the relation between the number of e-foldings N and model parameters. We are interested in the moment when the perturbation modes of conformal momentum

k , corresponding to the WMAP pivot scale $k/a_0 = 0.002 \text{ Mpc}^{-1}$ [12], exit the horizon at inflationary stage. Hereafter a_0 is the present scale factor, “ e ” and asterisk in subscript refer to the values at the end of inflation and at the horizon crossing of the modes defined above; so $N = \log(a_e/a_*)$ by definition. At horizon crossing the Hubble parameter \mathcal{H} obeys the equation

$$\mathcal{H}_* = \frac{k}{a_*} \equiv \frac{k}{a_e} e^N.$$

Then

$$N = \log \frac{a_0 \mathcal{H}_*}{k} + \log \frac{a_r}{a_0} + \log \frac{a_e}{a_r}, \quad (3.23)$$

where “ r ” in subscript refers to the values at reheating. At the hot stage entropy in comoving volume is conserved, hence

$$\frac{a_r}{a_0} = \frac{T_0}{T_r} \frac{g_0^{1/3}}{g_r^{1/3}}, \quad (3.24)$$

where for the effective number of the SM degrees of freedom at present time and at reheating one has $g_0 = 43/11$ and $g_r = 106.75$, respectively. Since between inflation and reheating the Universe in our model expands as it would be at radiation domination, we have

$$\frac{a_e}{a_r} = \frac{T_r}{U_e^{1/4}} \frac{g_r^{1/4} \pi^{1/2}}{30^{1/4}}. \quad (3.25)$$

Plugging eqs. (3.24), (3.25) into (3.23) and substituting the Hubble parameter from the Friedman equation we obtain

$$N = \log \frac{T_0 g_0^{1/3}}{k/a_0} + \frac{1}{4} \log \frac{\pi^2}{270 g_r^{1/3}} + \log \frac{U_*^{1/2}}{M_P U_e^{1/4}}, \quad (3.26)$$

where T_0 is CMB temperature at present. Note that dependence of N on the reheating temperature enters only through the number of degrees of freedom at reheating g_r . It is due to the fact that at hot stage the expansion is not isothermal but an adiabatic process, so the cosmological observables depends on the effective numbers of degrees of freedom in plasma at the onset of the hot stage g_r and at present g_0 .

Numerically, for WMAP pivot scale $k/a_0 = 0.002 \text{ Mpc}^{-1}$ [12], one gets

$$N = 64.3 + \log \frac{U_*^{1/4}/M_P}{\epsilon_*^{1/4}} + \log \frac{U_*^{1/4} \epsilon_*^{1/4}}{U_e^{1/4}}. \quad (3.27)$$

The first logarithmic term in (3.27) follows from eq. (3.8), $\log \frac{U_*^{1/4}/M_P}{\epsilon_*^{1/4}} = -3.6$. Substituting eqs. (3.4), (3.2), (3.6) into the last logarithmic term in (3.27), one arrives at

$$N = 60.7 + \frac{1}{2} \log \frac{(1 + 16\xi + \sqrt{1 + 32(1 + 6\xi)\xi}) X_N/M_P}{8\sqrt{2} \sqrt{1 + (1 + 6\xi)\xi} X_N^2/M_P^2 (1 + \xi X_N^2/M_P^2)}. \quad (3.28)$$

We use this formula for the numerical estimates of the number of e-foldings in a model with non-zero finite ξ . One gets $N = 61.5$ – 58.5 for ξ changing from zero to infinity.

Let us analyze briefly the reheating itself. After the inflation the action for the scalar sector in the Einstein frame is

$$S = \int d^4x \left\{ \frac{M_P^2}{2} R + \frac{\partial_\mu H^\dagger \partial_\mu H}{\Omega^2} + \frac{\Omega^2 + 6\xi^2 X^2/M_P^2}{\Omega^4} \frac{\partial_\mu X \partial_\mu X}{2} - \frac{V(H, X)}{\Omega^4} \right\}.$$

In the reheating regime $X < X_e$ for $\xi < 1$ we can neglect the coefficient in front of the kinetic term for X : the Universe is dominated by the field X oscillating due to the quartic potential. For zero ξ the reheating due to the potential mixing term $\alpha X^2 H^\dagger H$ was analyzed in detail in [27], and leads to the bound

$$\alpha > 0.7 \times 10^{-11}.$$

For $\xi \neq 0$ additional contributions to the scattering process appear suppressed² by at least M_P^2/ξ . It is easy to estimate that this contribution is negligible compared to the quartic mixing term and all the bounds on α in [27] remain valid.

4. Phenomenology: searches in particle physics experiments

We see that models with very small non-minimal coupling $\xi < 2 \times 10^{-3}$ are disfavored from cosmology (upper panels in Fig. 1) and hence the relevant for low energy phenomenology quartic self-coupling β should obey $\beta > 3 \times 10^{-13}$ (lower right panel in Fig. 1). Accounting for ξ -dependence of β in eq. (2.3), where the SM Higgs boson mass is $m_h = 126$ GeV, and recalling that parameter α is bounded from above and below by constraints on quantum corrections and reheating temperature, respectively, as explained in Sec. 2, we end up with cosmologically motivated interval of the inflaton mass m_χ , confined between green and blue shaded regions of Fig. 2 at $\xi > 2 \times 10^{-3}$.

The upper limit on m_χ at a given ξ comes from requirement of sufficiently energetic reheating, that develops in the same way as in the case of $\xi = 0$, studied in Ref. [27]. It implies a limit $\alpha > 0.7 \times 10^{-11}$, which is substituted to (2.3) and with account of ξ -dependence of β is plotted in Fig. 2. The lower limit on m_χ for the case of $\xi = 0$ has been obtained in [1] from smallness of quantum corrections to the inflaton potential, guaranteed at

$$\alpha^2 < 0.1 \times \beta.$$

While $\xi < 1$, that estimate remains valid. We plug it into eq. (2.3) and recall ξ -dependence of β to depict the upper limit on inflaton mass in Fig. 2.

Another set of bounds on the model parameters comes from particle physics experiments. Thanks to the inflaton-Higgs mixing (2.4), inflaton can be produced in the same channels as the SM Higgs boson would be of the same mass m_χ . The same mixing is responsible for the inflaton decay to the light SM particles, and its decay pattern is exactly

²Let us note here the significant difference from the situation of $\xi \gg 1$, which is similar to the case of R^2 inflation [23, 26]. For $\xi \gg 1$ the kinetic term is significantly modified, and the field redefinition for X is needed during reheating, leading to a much weaker suppression of the interactions with the Higgs boson by first power of M_P .

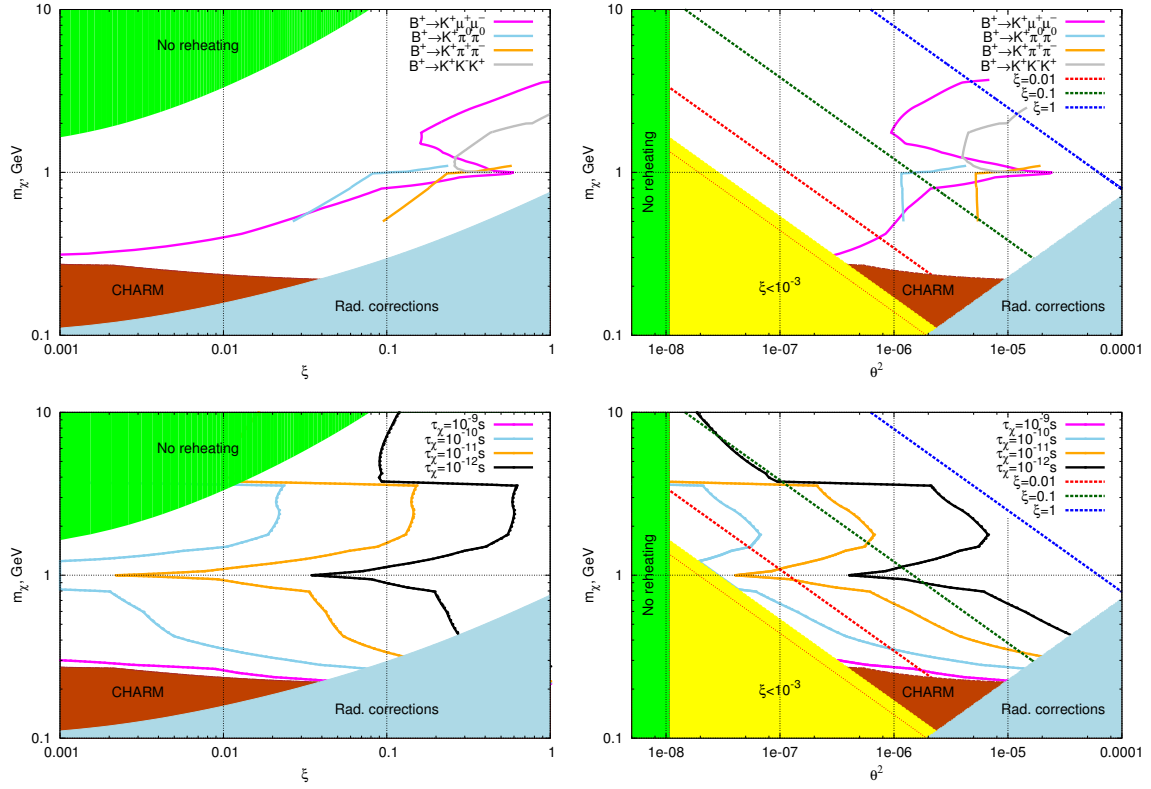


Figure 2: The bounds on the inflaton in the planes (m_χ, ξ) and (m_χ, θ^2) . Cosmological constraints are shaded: the green region leads to insufficient reheating, the blue region gives large radiative corrections to inflaton potential; the region $\xi < 2 \times 10^{-3}$ is disfavored from analysis of CMB data presented in Sec. 3. Other constraints are from direct searches in accelerator experiments: the brown region is excluded by the CHARM experiment. On the upper two plots other lines are the bounds from B meson decays, with the excluded areas to the right of the lines. On the lower plots the lines correspond to inflaton lifetime. Dashed lines on the right plots are isolines of constant ξ .

the same as that of the normal SM Higgs boson, if its mass would be equal to m_χ . Inflaton decay branching ratios are calculated in Ref. [1] and here we reproduce them in Fig. 3. Note that in the inflaton mass range 1-2 GeV the estimates are rather vague because of QCD-uncertainties with hadronization. The inflaton lifetime depends on β and hence on ξ given the relation in Fig. 1 (right lower panel). We have recalculated inflaton lifetime by replacing $\beta \rightarrow \beta(\xi)$ in formulas of Ref. [1], the result is plotted in Fig. 3 (right panel).

Inflaton-to-Higgs mixing opens a room for direct searches for the light inflaton, which has been thoroughly studied in Ref. [1]. Remarkably, only m_χ and β enter inflaton-Higgs mixing (2.4), hence no dependence on the Higgs boson mass in the inflaton production rate or its decay branching ratios. Thus laboratory experiments allow for probing directly the inflaton quartic coupling responsible for the inflationary stage in the early Universe. In [1] we have found as most relevant the analysis of the CHARM bound. In case of non-zero ξ it goes along the same lines as in Sec. 6 of Ref. [1], with the only substitution of $\beta \rightarrow \beta(\xi)$

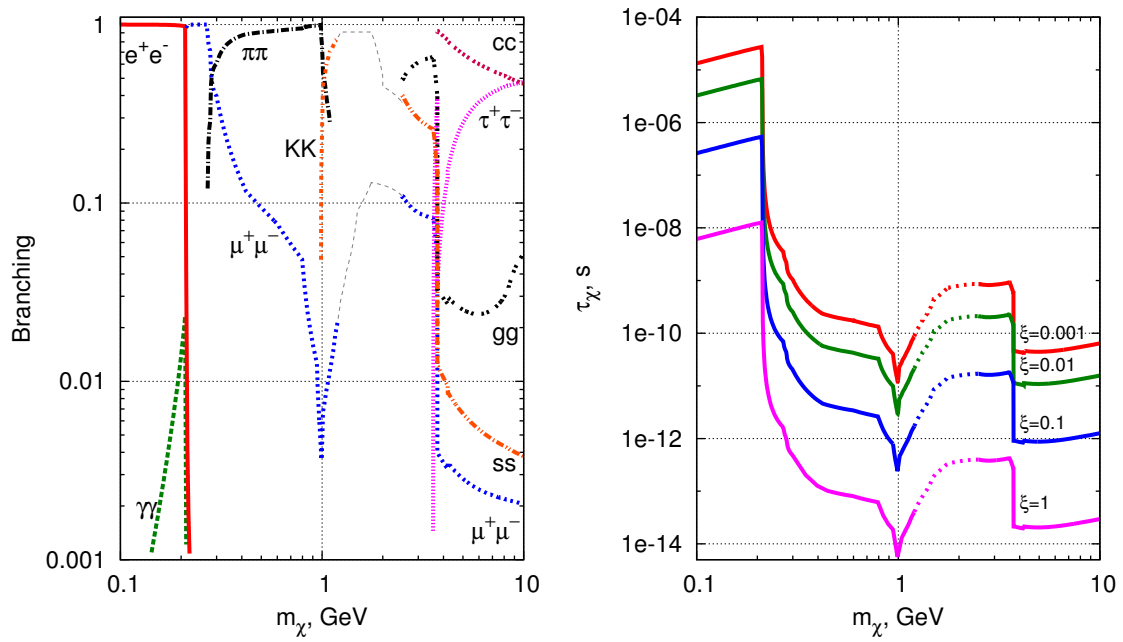


Figure 3: Inflaton decay branching ratios (*left plot*) and inflaton lifetime (*right plot*); theoretical predictions for $m_\chi \simeq 1 - 2 \text{ GeV}$ (thin dashed lines on the *left plot* and dotted lines on the *right plot*) suffer from significant QCD-uncertainties.

in the inflaton decay length and meson decay branching ratios into inflaton. The resulting bound on $\beta(\xi)$ is read from Fig. 5 of Ref. [1] and thus we obtain the lower limit on the inflaton mass at given ξ depicted in Fig. 2. One observes in Fig. 2, that with increasing ξ at first the lower masses become allowed, but quite fast the radiative corrections bound for the inflation starts to dominate. Moreover, the whole inflaton window moves to higher masses (that was in fact known, see [27]).

The most promising processes in searches for the inflaton of $m_\chi < 5 \text{ GeV}$ are b -quark decays. In Ref. [1] we found for the B -meson decay rate

$$\begin{aligned} \text{Br}(B \rightarrow \chi X_s) &\simeq 10^{-6} \times \left(1 - \frac{m_\chi^2}{m_b^2}\right)^2 \left(\frac{\beta(\xi)}{1.5 \times 10^{-13}}\right) \left(\frac{300 \text{ MeV}}{m_\chi}\right)^2 \\ &\simeq 4.8 \times 10^{-6} \times \left(1 - \frac{m_\chi^2}{m_b^2}\right)^2 \left(\frac{\theta^2}{10^{-6}}\right). \end{aligned} \quad (4.1)$$

Here X_s refers to a strange hadron, which most probably turns out to be a K -meson, because the other decay product, χ , is a scalar. For a not so small value of ξ the inflaton is short lived (see Fig. 3) and decays inside the detector into SM particles,

$$\chi \rightarrow \mu^+ \mu^-, \pi^+ \pi^-, \pi^0 \pi^0, K^+ K^-, \dots,$$

thus contributing to the corresponding three-body decay modes of B -meson. Conservatively, we require that inflaton contributions, estimated with eq. (4.1) and decay branching

ratios in Fig. 3, are smaller than the measured branching ratios of corresponding non-resonant three-body decays of B -mesons [28], specifically

$$\text{Br}(B \rightarrow K^+ \pi^0 \pi^0) < 1.6 \times 10^{-5}, \quad (4.2)$$

$$\text{Br}(B \rightarrow K^+ \pi^- \pi^+ \text{ nonresonant}) < 1.6 \times 10^{-5}, \quad (4.3)$$

$$\text{Br}(B \rightarrow K^+ \mu^+ \mu^-) < 4.8 \times 10^{-7}, \quad (4.4)$$

$$\text{Br}(B \rightarrow K^+ K^- K^+ \text{ nonresonant}) < 2.8 \times 10^{-5}. \quad (4.5)$$

Resulting limits on model parameters are presented in Fig. 2. One can see, that a wide range of inflaton masses $300 \text{ MeV} \lesssim m_\chi \lesssim 5 \text{ GeV}$ is allowed for direct investigation at B -factories. The branching ratio of the two-body B -meson decay is at the level of $10^{-5} - 10^{-8}$ for $\beta(\xi)$ in Fig. 1.

Note in passing that our limits from B -physics are only approximate. First, they do not apply to the inflaton masses close to the masses of known hadronic resonances, since limits (4.2)–(4.5) are obtained from the analysis of corresponding Dalitz plots with removed resonances. Second, the absence of new resonances on the Dalitz plots (would it be specially studied) put stronger constraints on the light inflaton model, than what we naively extracted from (4.2)–(4.5). Third, one should be careful about the bounds from $B \rightarrow K \mu \mu$, because it is not valid for living inflatons $\tau_\chi \gtrsim 10^{-9}$, which can escape the detector (the lower region near the CHARM bound).

5. Discussion

To summarize, following the new experimental results both in particle physics [8, 9] and cosmology [10, 12, 11], we have extended our previous analysis [1] of the large-field chaotic inflationary model with quartic inflaton potential to the case of inflaton non-minimally coupled to gravity. We concluded, that the range of cosmologically acceptable masses for the inflaton is wider, while the coupling between the inflaton and the SM sector can be stronger, than suggested [1] for the minimally coupled case. The model can be further tested with new cosmological data, especially with further results on the tensor perturbations of the CMB.

The best place to probe the model remains the same: studies of two-body B -meson decays to kaons and inflatons, where the inflaton decays further into electron, muon, pion or kaon pair. However with the increase of the non-minimal coupling ξ the inflaton lifetime decreases. In a smaller part of parameter space the inflaton lifetime is long enough to cover a macroscopic distance from the B -meson decay position, which leads to the following signatures of the inflaton event: either missing energy (if the inflaton escapes from the detector) or significantly displaced vertex of the inflaton decay. Therefore, searches for process $B \rightarrow K + \text{nothing}$ with *two-body kinematics* and $B \rightarrow K + \chi$ with subsequent decay of χ to a pair of SM light particles at a macroscopic distance from the B -meson decay are necessary to probe the model. In a wider region of parameter space the inflaton decays close to the B -meson decay point, thus effectively contributing to the three-body B -meson decays. This gives a possibility to probe the model by investigating neutral two-body resonance-like events ($\mu^+ \mu^-$, $\pi^+ \pi^-$, etc.) among three-body final states of B -meson

decays. The typical distance to the inflaton decay vertex, distinguishing these two options, can be deduced from the lifetime of the inflaton, see Figs. 2, 3.

Independently (and also if B -meson decays are kinematically forbidden) the model can be probed to some extent in a fixed-target experiment. The inflaton can be produced by protons on target either in meson decays or through the gluon fusion similar to the SM Higgs boson, would it be of the same mass. This option has been studied in detail in Ref. [1] for the case of the inflaton minimally-coupled to gravity. The estimates of the inflaton production cross sections for a realistic set of beams are given there in Fig. 4. For the non-minimally coupled inflaton those estimates have to be multiplied by a factor $\beta(\xi)/(1.5 \times 10^{-13})$, which scales from 1 to 3×10^3 when ξ grows from 0 to 1, see Fig. 1.

Note also, that the model depends on the positivity of the SM Higgs boson self-coupling for the whole range of energy scales, relevant for inflation (or, to be more precise, for the scales of the order $\sim \sqrt{\alpha/\lambda} M_P$). It implies a lower limit on the Higgs boson mass. This limit is slightly weaker than the bounds from the positivity of the Higgs self-coupling up to the Planck scale (see [29, 30, 31]), and is compatible with the observations at LHC within presently achieved precision for the Higgs and top-quark masses.

Finally, the inflationary model under discussion can be further modified (completed) to solve other phenomenological problems of the SM (neutrino oscillations, dark matter phenomena, baryon asymmetry of the Universe) in such a way that the inflationary dynamics and the low-energy phenomenology of the light inflaton remains unchanged. The working example is an extension of the SM with three Majorana sterile neutrinos, which can explain neutrino oscillations via type-I seesaw mechanism, the baryon asymmetry via resonant leptogenesis in the early Universe [32, 14] and the lightest of the sterile neutrinos can serve as dark matter. With sterile neutrinos in GeV mass range the model is known as ν MSM [13, 14] and thoroughly studied [33], though both seesaw mechanism and resonant leptogenesis allow for (much) heavier neutrinos as well. Note, that within our logic of only one energy scale in the whole model advertised in Introduction and Sec. 2, the natural mass scale for these neutrinos should not exceed the Electroweak scale.

Moreover, the inflaton being singlet with respect to the SM gauge group may have Yukawa coupling to the sterile neutrino and hence contributes to their production in the early Universe and their masses after getting non-zero vacuum expectation value. In this way one can avoid the explicit mass term for the sterile neutrino and deal with the explicitly one-scale model of particle physics. This particular setup has been considered in Appendix A of Ref. [1] and all formulas presented there can be fully transferred to the case of the non-minimally coupled to gravity inflaton, since all of them explicitly contain the factor $\beta = \beta(\xi)$ and the inflaton mass. Most importantly, in the non-minimally coupled case as well, the inflaton may contribute to the dark matter production (the lightest sterile neutrino) thus eluding the strong fine-tuning [34, 35] required in the simplest version of ν MSM to generate the dark matter. Remarkably, the Yukawa interactions of inflaton with sterile neutrinos is bounded from the requirement of small radiative corrections to the inflationary potential, and *can not give rise to very large masses* [1, 7] consistently with the logic adopted.

However, explicit mass parameters in the sterile neutrino part of the action are gener-

ally allowed. The only theoretical bound on them comes from the smallness of generated by the seesaw couplings one-loop finite corrections to the SM Higgs boson mass. The bound allows the sterile neutrino masses to be well above the EW scale. In this setup the mechanism of generation of the dark matter sterile neutrino in the inflaton decays still works, as well as the resonant leptogenesis [32]. We do not study this setup further as one beyond our logical framework.

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